

The Generalized Dilaton Supersymmetry Breaking Scenario ^{*}

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Abstract

We show that the usual dilaton dominance scenario, derived from the tree level Kähler potential, can never correspond to a global minimum of the potential at $V = 0$. Similarly, under very general assumptions it cannot correspond to a local minimum either, unless a really big conspiracy of different contributions to the superpotential $W(S)$ takes place. These results, plus the fact that the Kähler potential is likely to receive sizeable string non-perturbative contributions, strongly suggest to consider a more general scenario, leaving the Kähler potential arbitrary. In this way we obtain generalized expressions for the soft breaking terms but a predictive scenario still arises. Finally, we explore the phenomenological capability of some theoretically motivated forms for non-perturbative Kähler potentials, showing that it is easy to stabilize the dilaton at the realistic value $S \sim 2$ with just one condensate and no fine-tuning.

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1 Introduction and outline

The so-called “Dilaton Dominance” [1, 2] model has become a popular scenario of supersymmetry (SUSY) breaking in the framework of superstrings. The basic assumptions of this scenario are the following.

1. Supersymmetry is broken in the dilaton field (S) sector. In other words, only the F_S auxiliary field is to take a non-vanishing VEV.
2. The dilaton dependence of the Kähler potential, K , is assumed to be sufficiently well approximated by the tree-level expression, $K = -\log(S + \bar{S})$.
3. The superpotential W is in such an (unknown) way that the minimum of the potential lies at an acceptable value for the dilaton. More precisely, since $\langle \text{Re}S \rangle = g_{string}^{-2}$, i.e. the unified gauge coupling constant at the string scale, $\langle S \rangle \simeq 2$ has to be assumed.
4. The condition of vanishing cosmological constant, i.e. $V = 0$ at the minimum of the potential is also normally assumed.

The first assumption is plausible since the dilaton field is always present in any string scenario (at least in the weak coupling approach) and it has Planck-mass suppressed couplings with all the matter fields, i.e. it is a “hidden-sector” field. The second assumption is attractive since the tree-level form of $K(S, \bar{S})$ is universal for any string construction, thus providing a model independent framework. The third assumption is phenomenologically mandatory, while the last assumption is the usual one in supergravity (SUGRA) models¹.

The previous assumptions lead to some interesting relationships among the different soft terms, with an automatical implementation of universality, something which is phenomenologically welcome for FCNC reasons [4].

The interest of this scenario raises two questions:

- i)* Is there any form of the superpotential $W(S)$ (preferably with some theoretical justification) able to fulfill the previous requirements (i.e. a minimum of the scalar potential at $\langle \text{Re}S \rangle \simeq 2$ and $V = 0$)?
- ii)* How good is expected to be the tree-level approximation for $K(S, \bar{S})$?

Regarding the first question, (*i*), we will present in sect.2 a simple analytical argument that proves that, even if the potential has a minimum at $\langle \text{Re}S \rangle \simeq 2$ and $V = 0$, there must exist an additional minimum (or unbounded from below direction) in the perturbative region of S -values for which $V < 0$. This can be proven for an arbitrary form $W(S)$. Similarly, we will show that the very existence of such a (local) minimum is forbidden unless a really huge conspiracy of different contributions to $W(S)$ takes place.

¹For a discussion of the effect of the radiative corrections on the cosmological constant, see e.g. ref.[3].

Concerning the second question, (ii), there are, unfortunately, indications that the stringy non-perturbative corrections to the Kähler potential may be sizeable [5, 6]. In principle, this may be good news since it could help to avoid the above-mentioned problems. The trouble is that very little is known about the form of these non-perturbative corrections. On the other hand, there are string arguments [5] indicating that the stringy perturbative and non-perturbative corrections to W (and to the gauge kinetic function, f , beyond one-loop) are negligible, so $W(S)$ must be dominated by field-theory non-perturbative corrections (in particular gaugino condensation), whose general form is known. In any case, symmetries plus analyticity arguments strongly constrain [5] the functional dependence of $W(S)$, which must be $\sim \sum_i d_i e^{-a_i S}$.

From the previous reasons, it is suggestive to explore a generalized dilaton-dominated scenario in which $K(S, \bar{S})$ is left as an *arbitrary* function, while $W(S)$ is assumed to be as above. In some sense, this precisely represents a philosophy opposite to the “standard” (tree-level) dilaton-dominated scenario, but in our opinion much more likely to be close to the actual facts. In particular, it is interesting to analyze whether it is possible in this framework to avoid the shortcomings which are present in the tree-level one. This scenario is investigated in sect.3, where generalized expressions for the soft breaking terms are obtained.

Finally, it is tempting to examine the phenomenological capabilities of specific forms for the stringy non-perturbative contributions to K suggested by theoretical work in the subject [6]. This will be done in sect.4

2 Problems for the dilaton-dominated scenario

Taking the tree-level expression for the Kähler potential

$$K = -\log(S + \bar{S}) + \hat{K}(T, \bar{T}, \phi_I, \bar{\phi}_I) \quad , \quad (1)$$

where T, ϕ_I denote generically all the moduli and matter fields respectively, the scalar potential in the dilaton-dominance assumption reads

$$V = \frac{1}{2 \operatorname{Re} S} \left\{ |(2 \operatorname{Re} S) W_S - W|^2 - 3 |W|^2 \right\} \quad (2)$$

with $W_S \equiv \partial W / \partial S$. If the scenario is realistic, the previous potential should have a minimum at a realistic value of S , say S_0 , with

$$\operatorname{Re} S_0 = \frac{1}{g^2} \simeq 2 \quad , \quad (3)$$

where, for simplicity of notation, g denotes the gauge coupling constant at the string scale ($\sim 10^{17}$ GeV). In addition, the cosmological constant should be vanishing, i.e. $V(S_0) = 0$. This implies

$$|(2 \operatorname{Re} S_0) W_S(S_0) - W(S_0)| = \sqrt{3} |W(S_0)| \quad (4)$$

and, thus

$$\frac{2 \operatorname{Re} S_0}{\sqrt{3} + 1} \leq \left| \frac{W}{W_S} \right|_{S=S_0} \leq \frac{2 \operatorname{Re} S_0}{\sqrt{3} - 1} \quad (5)$$

Performing the following change of variables

$$z = e^{-\beta S} \quad (\beta \text{ arbitrary}) \quad (6)$$

the physical region of S , i.e. $\text{Re}S > 0$, is mapped into the circle of radius 1 in the z -plane. The “realistic minimum” point, $z_0 = e^{-\beta S_0}$, lies somewhere inside the circle. In the new variable, the functions W, W_S are written as

$$\begin{aligned} W(S) &= \Omega(z) \\ W_S(S) &= \Omega_S(z) \equiv -\beta z \Omega'(z) \end{aligned} \quad (7)$$

and condition (5) becomes

$$\frac{2 \text{Re}S_0}{\sqrt{3} + 1} \leq \left| \frac{\Omega(z)}{\Omega_S(z)} \right|_{z=z_0} \leq \frac{2 \text{Re}S_0}{\sqrt{3} - 1} \quad (8)$$

with $\text{Re}S_0 = -\log |z_0|/\beta$. Let us consider now the function

$$\rho(z) \equiv \frac{\Omega(z)}{\Omega_S(z)}. \quad (9)$$

Let us suppose for the moment that $\rho(z)$ is an analytical function with no poles inside the physical region $|z| < 1$. Then, the maximum of $|\rho(z)|$ in the region $|z| \leq |z_0|$ must necessarily occur (principle of maximum) at some point z_M belonging to the boundary, namely the circle $\mathcal{C} \equiv \{|z| = |z_0|\}$. If we consider the larger region enclosed by the broader circle $\mathcal{C}' \equiv \{|z| = |z'_0|\}$, with $|z'_0| > |z_0|$, the new maximum of $|\rho(z)|$ must occur now at some point, say $z = z_1$, belonging to the boundary \mathcal{C}' . From (8) it is clear that at z_1

$$|\rho(z_1)| = \left| \frac{\Omega(z)}{\Omega_S(z)} \right|_{z=z_1} > \frac{2 \text{Re}S_0}{\sqrt{3} + 1} \quad (10)$$

Taking the radius of \mathcal{C}' so that $\text{Re}S_1 \equiv -\log |z_1|/\beta = \frac{\sqrt{3}-1}{\sqrt{3}+1} \text{Re}S_0$, we can write (10) as

$$\left| \frac{\Omega(z)}{\Omega_S(z)} \right|_{z=z_1} = \left| \frac{W}{W_S} \right|_{S=S_1} > \frac{2 \text{Re}S_1}{\sqrt{3} - 1} \quad (11)$$

Therefore, at $S = S_1$ the potential (2) has a negative value. On the other hand S_1 still belongs to the perturbative region

$$\text{Re}S_1 \simeq 0.27 \text{Re}S_0 \quad (12)$$

which means $\alpha \simeq 0.14$. The only way-out to the previous argument is to allow the function $\rho(z) \equiv \frac{\Omega(z)}{\Omega_S(z)}$ to have some pole in the region enclosed by \mathcal{C}' . But then $\left| \frac{\Omega}{\Omega_S} \right| \rightarrow \infty$ near the pole and necessarily $\left| \frac{\Omega}{\Omega_S} \right| > \frac{2 \text{Re}S}{\sqrt{3}-1}$ at some point with non-zero Ω . Hence we arrive to the same conclusion.

The previous argument shows that the realistic minimum assumed to take place in the usual dilaton-dominated scenario can *never* correspond to a global minimum. This is not certainly the most desirable situation.

We can go a bit further and show that under very general assumptions, the realistic point $S = S_0$ cannot correspond to $V(S_0) = 0$.

From symmetry and analyticity arguments we know [5] that the non-perturbative superpotential must take the form $W = \sum_i d_i e^{-a_i S}$, so it is reasonable to assume that at the realistic point ($S_0 \sim 2$) W is dominated by one of the terms, say $W \sim e^{-aS}$ (as it happens for instance in usual gaugino condensation). Then, the vanishing of Λ_{cos} at $S = S_0$, eq.(4), implies

$$\left(-a - \frac{1}{2\text{Re}S_0}\right)^2 - \frac{3}{4(\text{Re}S_0)^2} = 0. \quad (13)$$

From (13) we obtain $a \simeq (\sqrt{3} - 1)/4$, absolutely incompatible with the requirement of a hierarchically small SUSY breaking (note that $\langle W \rangle \sim 1$ TeV requires $a \simeq 18$).

It could happen, however, that two or more terms of the form $W = \sum d_i e^{-a_i S}$ cooperate at the particular region $S \simeq 2$ to produce a more realistic scenario². It is interesting to show that for this to happen in a dilaton dominated scenario a really huge conspiracy must take place. From eq.(5) we see that the condition $V(S_0) = 0$ implies

$$|W_S| \sim |W| \quad (14)$$

Since $W_S = -\sum a_i d_i e^{-a_i S}$ and the condition of hierarchically SUSY breaking requires $a_i \gtrsim O(10)$, it is clear that a cancellation between terms with different exponents must occur inside W_S for (14) to be fulfilled. This can happen if the d_i coefficients are also (very) different, but is not enough: The extremalization condition $\partial V/\partial S = 0$ applied to the potential (2) leads to two possible solutions,

$$\begin{aligned} 2 \text{Re}S W_S - W &= 0 \\ (2 \text{Re}S)^2 W_{SS} &= 2W^* \frac{2 \text{Re}S W_S - W}{(2 \text{Re}S W_S - W)^*}. \end{aligned} \quad (15)$$

Since $(2 \text{Re}S W_S - W)$ is proportional to F_S , only the second solution in (15) is compatible with a dilaton dominance scenario³. This requires

$$(2 \text{Re}S)^2 |W_{SS}| = 2|W| \quad (16)$$

This is an unnatural requirement since the typical size of W_{SS} is $W_{SS} \sim a_i^2 W \gg W$. So, a second unpleasant cancellation must occur at $S = S_0$. Finally, one must demand that the extremum corresponds to a minimum. The determinant of the Hessian matrix (evaluated in the variables $\text{Re}S, \text{Im}S$) reads

$$\mathcal{H} = -4 \left\{ (2 \text{Re}S)^2 |3W_{SS}W_S + W_{SSS}(2 \text{Re}S W_S - W)^*|^2 - \frac{16}{(2 \text{Re}S)^4} |W|^4 \right\} \quad (17)$$

²This is the mechanism of the so-called racetrack models to generate SUSY breaking (see e.g. [7] and references therein). These models, however, lead naturally to moduli dominance SUSY breaking rather than dilaton dominance. For an attempt to generate dilaton dominance in this way see ref.[8]

³The first solution of eq.(15) was shown in ref.[9] to be incompatible with a global minimum at $V = 0$, unless a certain number of terms in the expansion $W = \sum d_i e^{-a_i S}$ were relevant, in concordance with the general result obtained between eqs.(4)–(12).

Using eqs.(4, 14, 16), the condition for a minimum, $\mathcal{H} > 0$, requires

$$\frac{2}{\sqrt{3}}(\text{Re}S)^2|W_{SSS}| \sim |W_S| \sim |W|. \quad (18)$$

Again, this is an (even more) unnatural requirement, since typically $|W_{SSS}| \sim a_i^3|W|$, so a third fine cancellation must take place.

To summarize the results of this section, the standard (tree-level) dilaton-dominated scenario can *never* correspond to a global minimum of the potential at $V = 0$. Similarly, under very general assumptions it cannot correspond to a local minimum either, unless a really big conspiracy of different contributions to $W(S)$ takes place. In addition, let us mention that it has recently been shown [10] that the effective low-energy scenario to which it gives place necessarily contains dangerous charge and color breaking minima.

3 The generalized dilaton-dominated scenario

The previous results, plus the fact that the Kähler potential is likely to receive sizeable string non-perturbative corrections, strongly suggest to consider a more general scenario, as commented in the introduction. Thus, in this section we will study how far we can go with the usual assumption of “dilaton–dominance” (i.e. only $|F_S| \neq 0$), but leaving the Kähler potential arbitrary. The potential reads

$$V = K_{S\bar{S}} |F_S|^2 - 3e^K |W|^2, \quad (19)$$

where

$$F_S = e^{K/2} \left\{ \left(K^{-1} \right)^{\bar{\phi}S} (\partial_{\phi} W + W K_{\phi})^* \right\} \quad (20)$$

with ϕ running over all the chiral fields. If we also assume vanishing cosmological constant, then

$$|F_S| (K_{S\bar{S}})^{1/2} = \sqrt{3} e^{K/2} |W| = \sqrt{3} m_{3/2}. \quad (21)$$

The passage to the effective low-energy theory involves a number of rescalings. In particular, the canonically normalized scalar and gaugino fields are given by $\hat{\phi}_I = K_{I\bar{I}} \phi_I$, $\hat{\lambda}_a = (\text{Re} f_a)^{1/2} \lambda_a$ respectively, where $K_{I\bar{I}} = \partial_{\phi_I} \partial_{\bar{\phi}_I} K$ and f_a is the gauge kinetic function. Likewise, the Yukawa couplings of the effective superpotential are given by $\hat{Y}_{IJL} = e^{K/2} (W^*/|W|) (K_{I\bar{I}} K_{J\bar{J}} K_{L\bar{L}})^{-1/2} Y_{IJL}$. With these redefinitions the supersymmetric part of the Lagrangian is obtained by applying the usual global-SUSY rules, while the soft part of the Lagrangian is given by

$$- \mathcal{L}_{\text{soft}} = \frac{1}{2} M_{1/2}^a \hat{\lambda}_a \hat{\lambda}_a + \sum_I m_I^2 |\hat{\phi}_I|^2 + \left(A_{IJL} Y_{IJL} \hat{\phi}_I \hat{\phi}_J \hat{\phi}_L + \text{h.c.} \right) + \dots \quad (22)$$

The values of the gaugino masses, $M_{1/2}^a$, scalar masses, m_I^2 , and coefficients of the trilinear scalar terms, A_{IJL} , can be computed using general formulae [1, 11] and eqs.(19–21)

$$|M_{1/2}| = \frac{1}{2} \sqrt{3} g^2 m_{3/2} (K_{S\bar{S}})^{-1/2} \quad (23)$$

$$m_I^2 = m_{3/2}^2 \left[1 - 3K_{S\bar{S}} \left(\frac{\partial_S \partial_{\bar{S}} K_{I\bar{I}}}{K_{I\bar{I}}} - \frac{|\partial_S K_{I\bar{I}}|^2}{K_{I\bar{I}}^2} \right) \right] \quad (24)$$

$$|A_{IJL}| = \sqrt{3} m_{3/2} (K_{S\bar{S}})^{-1/2} \left| K_S + \sum_{p=I,J,L} \left(\frac{\partial_S K_{p\bar{p}}}{K_{p\bar{p}}} \right) \right|. \quad (25)$$

We have assumed here that the Yukawa couplings appearing in the original superpotential, W , do not depend on the dilaton S . This is true at tree-level (and thus at the perturbative level) and, since they are parameters of the superpotential, they are not likely to be appreciably changed at the non-perturbative level [5]. The expression for the coefficient of the bilinear term, B , depends on the mechanism of generation of the μ term, so we prefer to leave it as an independent parameter. In the previous equations, besides the value of $m_{3/2}$, there are four unknowns: $K_S, K_{S\bar{S}}, \frac{\partial_S K_{I\bar{I}}}{K_{I\bar{I}}}, \frac{\partial_S \partial_{\bar{S}} K_{I\bar{I}}}{K_{I\bar{I}}}$. From the phenomenological requirement of (approximate) universality [4] the quantities $\frac{\partial_S K_{I\bar{I}}}{K_{I\bar{I}}}, \frac{\partial_S \partial_{\bar{S}} K_{I\bar{I}}}{K_{I\bar{I}}}$ should be universal for all the ϕ_I (matter) fields. If we further assume that $K_{I\bar{I}}$ does not get S -dependent contributions⁴ then there are just two unknowns: $K_S, K_{S\bar{S}}$ (besides the value of $m_{3/2}$). We obtain the following relations:

$$\begin{aligned} m_I^2 &= m_{3/2}^2 \\ \left| \frac{M_{1/2}}{m_{3/2}} \right| &= \left[\frac{3g^4}{4K_{S\bar{S}}} \right]^{1/2} \\ \left| \frac{M_{1/2}}{A} \right| &= \left| \frac{g^2}{2K_S} \right|. \end{aligned} \quad (26)$$

Notice that in the tree level limit $K_S = \frac{1}{2\text{Re}S}$, $K_{S\bar{S}} = \frac{1}{4(\text{Re}S)^2}$, and we recover from eqs.(26) the usual tree level relations [1, 2]

Now, the non-perturbative superpotential must take the form $W = \sum_i d_i e^{-a_i S}$, so, as explained above, it is reasonable to assume that at the realistic point, $S \sim 2$, W is dominated by one of the terms, say $W \sim e^{-aS}$, as it happens for instance in usual gaugino condensation. Thus, the condition $\Lambda_{\cos} = 0$, i.e. eq.(21), gives us a further constraint, namely

$$|-a + K_S|^2 - 3K_{S\bar{S}} = 0 \quad (27)$$

which relates the values of $K_S, K_{S\bar{S}}$ and a . Notice that the latter is essentially fixed by the condition of a hierarchical SUSY breaking ($a \simeq 18$), thus eqs.(26) and eq.(27) give a non-trivial scenario whose phenomenology could be investigated.

In obtaining eq.(27) from (21) and (20) we have ignored the possible dependence of W and K on the T -field (or fields) as well as the mixing between S and T . This is justified by the following. Working, for the sake of simplicity, with the dilaton S and an overall modulus T the explicit form of F_S is given by

$$F_S = e^{K/2} \left\{ \left(K^{-1} \right)^{\bar{S}S} (\partial_S W + W K_S)^* + \left(K^{-1} \right)^{\bar{T}S} (\partial_T W + W K_T)^* \right\} \quad (28)$$

⁴ This, of course, may not occur. However, it is a common assumption of all existing string-based models (including the usual dilaton-dominance model). There are some examples (mainly orbifold compactifications) where the one-loop corrections to $K_{I\bar{I}}$ are known [12] and have been incorporated in the analysis [13, 2], but they are very small.

The fact that, by assumption, $F_T = 0$ at the minimum of the potential, implies $(\partial_T W + W K_T)^* = \frac{(K^{-1})^{\bar{S}T}}{(K^{-1})^{TT}} (\partial_S W + W K_S)^*$. Thus $F_S = e^{K/2} [(K^{-1})^{\bar{T}T}]^{-1} |K^{-1}| (\partial_S W + W K_S)^* = e^{K/2} (K_{\bar{S}S})^{-1} (\partial_S W + W K_S)^*$, and the scalar potential (19) reads

$$V = e^K \left\{ (K_{\bar{S}S})^{-1} |\partial_S W + W K_S|^2 - 3 |W|^2 \right\} \quad (29)$$

where the explicit dependence on T is irrelevant. Then, using $W \simeq e^{-aS}$ the condition $V = 0$ translates into eq.(27) above.

4 Ansatz for non-perturbative Kähler potentials

It is tempting to go a bit further in our analysis and explore the phenomenological capabilities of explicit (stringy) non-perturbative effects on K which have been suggested in the literature⁵.

From the arguments explained in ref.[5], it is enough to focuss our attention on possible forms of $K(S+\bar{S})$, exploring the chances of getting a global minimum at $V = 0$. From eq.(27) we can check that a successful generalized dilaton dominated scenario (i.e. with $V = 0$ at $\text{Re}S \simeq 2$) requires very sizeable non-perturbative corrections to K . This can be seen by considering the form of eq.(27) when K is substituted by the tree level expression $K = -\log(S + \bar{S})$. Then, the left hand side of (27) reads

$$\left[\left(-a - \frac{1}{2\text{Re}S} \right)^2 - \frac{3}{4(\text{Re}S)^2} \right]_{\text{Re}S \simeq 2}, \quad (30)$$

which, as mentioned in sect.2, cannot be cancelled for any reasonable value of a . The fact that $a \simeq 18$ implies that the non-perturbative “corrections” to $K_S, K_{\bar{S}S}$ must indeed be bigger than the tree level values. If we further demand that the potential has a minimum at $V = 0$ this implies

$$\begin{aligned} \left(-a + \frac{1}{2}K' \right) K'' - \frac{3}{4}K''' &= 0 \\ \left(-a + \frac{1}{2}K' \right)^2 - \frac{3}{4}K'' &= 0 \end{aligned} \quad (31)$$

(at the realistic point $\text{Re}S = 2$), where the primes denote derivatives with respect to $\text{Re}S$ (the second equation is simply eq.(27)). It would be nice if the previous conditions could be fulfilled by some simple form of K .

Since in the known examples (mainly orbifold compactifications) the perturbative corrections to K are small, we can tentatively write $K(\text{Re}S)$ as

$$K(\text{Re}S) = -\log(2 \text{Re}S) + K_{np}(\text{Re}S) \quad (32)$$

⁵For other attempts in this sense, see ref.[14].

where the first term corresponds to the tree level expression and K_{np} denotes the non-perturbative contributions.

The next step is to choose some plausible form for $K_{np}(\text{Re}S)$ and to study the form of V , see eq.(29). A reasonable assumption could be to impose that $K_{np}(\text{Re}S) \rightarrow 0$ for $\text{Re}S \rightarrow \infty$ (i.e. the limit of vanishing gauge coupling) and that $K_{np}(\text{Re}S)$ is zero at all order in the perturbative expansion. The field theory contributions to K_{np} have been evaluated in ref.[15], but for a realistic case they turn out to be too tiny to appreciably modify the tree-level results⁶. On the other hand, according to the work of ref.[6], stringy non-perturbative effects may be sizeable (even for weak four-dimensional gauge coupling) and plausibly go as $g^{-p}e^{-b/g}$ with $p, b \sim O(1)$. Then a simple possibility is to take

$$K_{np} = d g^{-p} e^{-b/g} \quad (33)$$

where d, p, b are constants with $p, b > 0$ and $g^{-2} = \text{Re}S$. Then, it can be explicitly shown that for $p = 1$ the conditions (31) can be satisfied for $\text{Re}S = 2$ and $b \simeq 1$, leading indeed to the appearance of a minimum with $V = 0$. Unfortunately, the required value of d is very large (and negative), which means that the model is not realistic, since $m_{3/2}^2 = e^K |W|^2$ would be extremely small.

Perhaps a more sensible approach is to write the decomposition (32) for the exponential of the Kähler potential, i.e.

$$e^K = e^{K_{tree}} + e^{K_{np}}, \quad (34)$$

as is suggested by the fact that the SUGRA lagrangian has an exponential dependence on K (e.g. it can be written as $\mathcal{L} = [e^{-K/3}]_D$), and then take $e^{K_{np}} \sim d g^{-p} e^{-b/g}$. (For small K_{np} eqs. (32) and (34) would be essentially equivalent.)

The numerical results indicate that, plugging this ansatz, it is not difficult to get a minimum of the potential using just one condensate with $a \simeq 18$ (thus guaranteeing the correct size of SUSY breaking) and sensible values for the d, p, b constants. More precisely, for $d = -3$, $p = 0$, $b = 1$, there is a minimum of the potential at $\text{Re}S = 2.1$. Unfortunately, the negative value of d makes this example unacceptable. Another example with positive d is $d = 7.8$, $p = 1$, $b = 1$, which has a minimum at $\text{Re}S = 1.8$. However, as a general result, playing just with these simple forms for K , it seems impossible to get the minimum at $V = 0$. Anyway, it is impressive that just with one condensate and choosing very reasonable values for d , p and b (note also that there is no fine-tuning in the previous choices), a minimum appears at the right value of the dilaton. Another typical characteristic of these examples is the appearance of singularities in the potential caused by zeroes of the second derivative of the Kähler potential, K'' , at some particular values of S . Of course, this can be cured by additional terms in K .

Finally, it would be worth to analyze the possible implications of this kind of non-perturbative Kähler potentials for other different matters, such as the cosmological moduli problem [16, 17]

⁶ More precisely, the authors of ref.[15] obtain $K = -3 \log [(2 \text{Re}S)^{1/3} + e^{-K_p/3} + d e^{-(2 \text{Re}S)/2c}]$, where K_p denotes the perturbative corrections to the tree-level expression, d is an unknown constant and c is essentially the coefficient of the one-loop beta function (the non-perturbative superpotential goes as $e^{-3S/2c}$).

5 Conclusions

We have shown that the usual dilaton dominance scenario, derived from the tree level Kähler potential, can *never* correspond to a global minimum of the potential at $V = 0$. Similarly, it cannot correspond to a local minimum either, unless a really big conspiracy of different contributions to the superpotential $W(S)$ takes place.

These results, plus the fact that the Kähler potential is likely to receive sizeable string non-perturbative corrections, strongly suggest to consider a more generalized scenario, leaving the Kähler potential arbitrary. In this way we obtain generalized expressions for the soft breaking terms. A predictive scenario arises if one further assumes, as it is usually done, that the matter fields kinetic terms, $K_{I\bar{I}}$, do not get sizeable S -dependent contributions and that the non-perturbative superpotential $W = \sum_i d_i e^{-a_i S}$ is dominated at the realistic point by one of the terms, say $W \simeq e^{-aS}$ (as it happens for instance in usual gaugino condensation). Then, the relevant formulae for the scalar and gaugino masses, $m_I^2, M_{1/2}$, and for the coefficient of the trilinear scalar terms, A , are given (see eqs.(26, 27)) by

$$\begin{aligned} m_I^2 &= m_{3/2}^2 \\ |M_{1/2}| &= m_{3/2} \frac{3g^2}{|K' - 2a|} \\ |A| &= |M_{1/2}| \left| \frac{g^2}{K'} \right|, \end{aligned} \tag{35}$$

where $K' \equiv \partial K / \partial (\text{Re} S)$ and $a \simeq 18$ from the condition of hierarchical SUSY breaking.

We have explored the phenomenological capability of some theoretically motivated forms of non-perturbative Kähler potentials, showing that for reasonable choices it is easy to get a minimum of the potential at the realistic value $S \sim 2$ with just one condensate and with no fine-tuning at all. The goal of getting a vanishing cosmological constant remains however as a more challenging task.

Finally, it would be nice if the new string-duality techniques (see in particular ref.[18]) could provide more precise information about the form of K , so its phenomenological viability could be tested in the sense exposed in this paper.

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